

A broad diphoton resonance at the TeV? Not alone

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The hint for a possible resonance in the diphoton channel with mass of 750 GeV disappeared in the data presented at ICHEP'16 by ATLAS and CMS. However, the diphoton final state remains as one of the golden channels for new physics discoveries at the TeV scale in the LHC experiments. This motivates us to analyze model-independently the implications of an $\mathcal{O}(\text{TeV})$ bump in the $\gamma\gamma$ final state. By means of forward sum-rules for $\gamma\gamma$ scattering, we show that a spin-0 resonance with mass of the order of the TeV and a sizable $\gamma\gamma$ partial width –of the order of a few GeV– must be accompanied by higher spin resonances with $J_R \geq 2$ with similar properties, as expected in strongly coupled extensions of the Standard Model or, alternatively, in higher dimensional deconstructed duals. Furthermore, independently of whether the putative $\mathcal{O}(\text{TeV})$ candidate is a scalar or a tensor, the large contribution to the forward sum-rules in the referred scenario implies the presence of states in the spectrum with $J_R \geq 2$, these high spin particles being a manifestation of new extra-dimensions or composite states of a new strong sector.

INTRODUCTION

Last December both ATLAS and CMS came up with significant local excesses in the diphoton spectrum at 750 GeV of more than 3σ and a global significance in the range of $1\text{--}2\sigma$ [1]. However, the results presented by both collaborations at the ICHEP'16 conference in Chicago [2] suggest that earlier data were a statistical fluctuation. Still, since the di-photon final state remains to be one of the most promising channels in new physics searches, we consider it might be interesting to point out some model-independent implications deriving from having a resonance with large $\gamma\gamma$ partial width, of the order of a few GeV.¹ Fully general quantum field-theory properties –analyticity, crossing symmetry and unitarity– imply a series of sum-rules for $\gamma\gamma$ scattering [4]. They require a precise cancellation between the different helicity contributions to the scattering amplitude that –remarkably– cannot be accomplished by just spin-0 intermediate exchanges. Likewise, a large resonance contribution cannot be cancelled by a weakly interacting background unless the non-resonant (non-R) loop diagrams get large and the theory abandons the perturbative regime. As a result of this, resonances with $J_R \geq 2$ must appear in this channel to fulfill the sum-rules. The exchange of states with spin $J_R \geq 2$ in the crossed channels leads to partial-wave amplitudes that diverge like $s^{J_R-1} \ln s$ (see e.g. [5]). This can only be solved through the infinite tower of resonances of increasing spin *à la Regge* that appears in strongly-coupled theories like Quantum Chromodynamics (QCD) or their equivalent duals –string theory scenarios *à la Veneziano* [6] or holographic models [7]–, these high spin particles being a manifestation of new extra-dimensions

or composite states of a new strong sector [8].

FORWARD SUM-RULES

We analyze the forward scattering amplitude ($t = 0$), $T_{\Delta\lambda}$, as a function of the variable $\nu \equiv (s - u)/2$,²

$$T_{\Delta\lambda}(\nu) = T(V(k, \lambda)V(k', \lambda') \rightarrow V(k, \lambda)V(k', \lambda')) , \quad (1)$$

between real spin-1 particles V with helicity difference $\Delta\lambda = |\lambda - \lambda'|$. In particular, we will focus on the case with $V = \gamma$, i.e., the forward $\gamma\gamma$ scattering.

Analyticity, crossing symmetry and unitarity imply the crossing and forward once-subtracted dispersion relation

$$T_{\Delta\lambda}(\nu) = T_{\overline{\Delta\lambda}}(-\nu) = T_{\Delta\lambda}(0) + \frac{\nu}{\pi} \int_{\nu_{\text{th}}}^{\infty} \frac{d\nu'}{\nu'} \left(\frac{\text{Im}T_{\Delta\lambda}(\nu' + i\epsilon)}{\nu' - \nu} - \frac{\text{Im}T_{\overline{\Delta\lambda}}(\nu' + i\epsilon)}{\nu' + \nu} \right) , \quad (2)$$

with the production threshold ν_{th} and $\overline{\Delta\lambda} \equiv 2 - \Delta\lambda$. The low-energy forward $\gamma\gamma$ scattering is provided by the Euler-Heisenberg effective field theory (EFT) [9],

$$T_{\Delta\lambda}(\nu) \stackrel{\nu \rightarrow 0}{\simeq} \mathcal{O}(\nu^2) . \quad (3)$$

This same result applies to other unbroken gauge theories (like, for instance, the scattering among same-color gluons, $g^a g^a \rightarrow g^a g^a$, that we do not discuss here). In the case of spontaneously broken theories it is possible to generate dimension 4 operators in the low-energy EFT, as in the forward ZZ scattering in the Standard Model

¹ See e.g. Ref. [3] for a review on weakly-coupled scenarios giving rise to a narrow-width resonance with mass in the TeV range.

² The use of the kinematical variable $\nu \equiv (s - u)/2$ is customary in fixed- t dispersive analyses of scattering amplitudes with definite $s \leftrightarrow u$ crossing properties (see in general Ref. [5]). In particular, for $t = 0$ one has $\nu = s$.

(SM) due to the tree-level Higgs (h) exchange [10]. We will focus in the following on the light-by-light scattering.

Matching Eq.(2) and the EFT result (3) up to $\mathcal{O}(\nu)$ leads to the forward photon sum-rule

$$0 = \frac{1}{\pi} \int_{\nu_{\text{th}}}^{\infty} \frac{d\nu'}{(\nu')^2} (\text{Im}T_2(\nu' + i\epsilon) - \text{Im}T_0(\nu' + i\epsilon)) , \quad (4)$$

equivalent to the Roy-Gerasimov-Moulin sum-rule for the inclusive $\gamma\gamma$ cross section $\sigma_{\Delta\lambda} = \sigma(\gamma(k, \lambda)\gamma(k', \lambda') \rightarrow X)$ [4]:

$$0 = \frac{1}{\pi} \int_{\nu_{\text{th}}}^{\infty} \frac{d\nu'}{\nu'} [\sigma_2(\nu') - \sigma_0(\nu')] , \quad (5)$$

by means of the relation $\sigma_{\Delta\lambda}(\nu') = \text{Im}T_{\Delta\lambda}(\nu' + i\epsilon)/\nu'$.

The resonant contribution to the spectral function (for $\nu \geq 0$) is given by

$$\begin{aligned} & \text{Im}T_{\Delta\lambda}(\nu + i\epsilon) \Big|_R \\ &= \sum_R 16\pi^2 (2J_R + 1) M_R \Gamma_{R \rightarrow [\gamma\gamma]_{\Delta\lambda}} \delta(\nu - M_R^2) , \end{aligned} \quad (6)$$

and turns (4) into

$$\begin{aligned} 0 = \sum_R 16\pi (2J_R + 1) \frac{(\Gamma_{R \rightarrow [\gamma\gamma]_2} - \Gamma_{R \rightarrow [\gamma\gamma]_0})}{M_R^3} \\ + \text{non-R}, \end{aligned} \quad (7)$$

where $\Gamma_{R \rightarrow [\gamma\gamma]_0}/2 = \Gamma_{R \rightarrow \gamma(+)\gamma(+)} = \Gamma_{R \rightarrow \gamma(-)\gamma(-)}$, $\Gamma_{R \rightarrow [\gamma\gamma]_2} = \Gamma_{R \rightarrow \gamma(+)\gamma(-)}$ and $\Gamma_{R \rightarrow \gamma\gamma} = \Gamma_{R \rightarrow [\gamma\gamma]_0} + \Gamma_{R \rightarrow [\gamma\gamma]_2}$. The ‘non-R’ term is provided by the contributions to the spectral function from loop diagrams in $\gamma\gamma \rightarrow \gamma\gamma$ without an intermediate s -channel resonance.

The importance of this sum-rule relies on the fact that the lowest-spin resonances ($J_R = 0$) only decay into $\gamma\gamma$ states with $\Delta\lambda = 0$ and hence give a negative contribution to the sum-rule (7). The positive terms with $\Gamma_{R \rightarrow [\gamma\gamma]_2}$ only appear for higher spin resonances with $J_R \geq 2$ [4, 11–13]:

$$\begin{aligned} \Gamma_{R \rightarrow \gamma\gamma} &= \Gamma_{R \rightarrow [\gamma\gamma]_0} , \quad \Gamma_{R \rightarrow [\gamma\gamma]_2} = 0 \quad \text{for } J_R = 0 , \\ \Gamma_{R \rightarrow \gamma\gamma} &= \Gamma_{R \rightarrow [\gamma\gamma]_2} , \quad \Gamma_{R \rightarrow [\gamma\gamma]_0} = 0 \quad \text{for } J_R = 3, 5, 7... \end{aligned} \quad (8)$$

Resonances with $J_R = 1$ are forbidden by the Landau-Yang theorem [14] and those with $J_R = 2, 4, 6...$ can in principle decay into $[\gamma\gamma]_0$ and $[\gamma\gamma]_2$ states [11]. In QCD the $\gamma\gamma$ decay of the lowest-lying spin-2 resonances ($\mathcal{T} = a_2, f_2, f'_2$) predominantly occurs with helicity $\Delta\lambda = 2$ [15], i.e., $\Gamma_{\mathcal{T} \rightarrow \gamma\gamma} \approx \Gamma_{\mathcal{T} \rightarrow [\gamma\gamma]_2}$. The sum-rule (4) is mostly saturated by the lightest pseudoscalar (π^0, η, η') and tensor (a_2, f_2, f'_2) mesons [12, 13], with the large spin-2 positive contribution cancelling out to a large extent the large negative spin-0 contribution. This situation resembles the case of (spin-2) massive gravitons G [7, 16, 17], where the decay $G \rightarrow V(\lambda)V(\lambda')$ always occurs with $\Delta\lambda = 2$ as the graviton couples to the stress-energy tensor of the gauge field $V = \gamma, g^a$.

Notice that the sum-rule studied here only relies on unitarity, analyticity and crossing symmetry and is therefore fulfilled in any possible beyond the SM (BSM) extension that assumes these general properties.

SUM-RULE WITH AN O(TEV) SPIN-0 RESONANCE

The hint for a possible diphoton resonance with mass 750 GeV in the 2015 [1] and early ($\sim 3fb^{-1}$) 2016 data [18] has not been confirmed analyzing the 15.4 (12.9) fb^{-1} data by the ATLAS (CMS) Collaboration. We wish, however, to illustrate how the sum rule Eq. (7) allows us to derive model independent constraints which may be useful to check the consistency of a supposed isolated future bump in the diphoton channel.

As our final purpose is to show the need of resonances with spin $J_R \geq 2$ in the spectrum, we will assume from now on that an O(TeV) candidate has spin zero; we will show that, even if it does not carry $J_R \geq 2$, the spectrum must contain further particles that do.

If the possible new state is a scalar (or a pseudoscalar), \mathcal{S} , the sum-rule (7) becomes

$$\begin{aligned} 16\pi \frac{\Gamma_{\mathcal{S} \rightarrow \gamma\gamma}}{M_{\mathcal{S}}^3} &= \sum_{R \neq \mathcal{S}} 16\pi (2J_R + 1) \frac{(\Gamma_{R \rightarrow [\gamma\gamma]_2} - \Gamma_{R \rightarrow [\gamma\gamma]_0})}{M_R^3} \\ &+ \text{non-R}. \end{aligned} \quad (9)$$

In order to fulfill the identity (9), resonances with $\Gamma_{R \rightarrow [\gamma\gamma]_2} \neq 0$ are needed on the right-hand side, i.e., resonances with spin $J_R \geq 2$. We also briefly discuss the importance of non-R loop contributions below.

In the case of a resonance with $M_{\mathcal{S}} \sim 1$ TeV (the errors are already very large at this energy to allow for this possibility) with a large $\gamma\gamma$ partial width (~ 10 GeV), the sum-rule (7) [19] gets the contribution

$$16\pi \frac{\Gamma_{\mathcal{S} \rightarrow \gamma\gamma}}{M_{\mathcal{S}}^3} \sim 0.5 \text{ TeV}^{-2} . \quad (10)$$

In comparison, the SM Higgs exchange yields the negligible contribution $16\pi\Gamma_{h \rightarrow \gamma\gamma}/m_h^3 \simeq 2.4 \times 10^{-4} \text{ TeV}^{-2}$ [20]. This large difference indicates that, in order to have a scalar resonance with a loop-induced decay like the SM one with similar scalar couplings, one needs either a large number of particles running in the intermediate loop or huge hypercharges [21].³

For radiatively generated $\mathcal{S} \rightarrow \gamma\gamma$ decays in weakly interacting theories in the TeV, small partial widths are expected. For a Higgs-like decay $\mathcal{S} \rightarrow \gamma\gamma$, one would have $\Gamma_{\mathcal{S} \rightarrow \gamma\gamma} \sim \alpha^2 M_{\mathcal{S}}^3 / (256\pi^3 v^2) |\sum_i N_i Q_i^2 A_i|^2$, summing

³ See also Ref. [22], where $\mathcal{O}(1)$ GeV width is advocated in the case of production via $\gamma\gamma$ fusion.

over the number N_i of particles i running within the loop with electric charge Q_i and the A_i coefficient from the one-loop integral (e.g., $|A_W| \leq 12.4$ and $|A_t| \leq 2.4$) [23]. This yields a contribution to (7) of the order of $16\pi\Gamma_{S\rightarrow\gamma\gamma}/M_S^3 \sim \alpha^2/(16\pi^2 v^2) \sim 10^{-5} \text{ TeV}^{-2}$.

Background non-R diagrams also contribute to the sum-rule Eq. (4). However, based on naive dimensional analysis, these are found to be small, of the order of

$$\frac{1}{\pi} \int_{\nu_{\text{th}}}^{\infty} \frac{d\nu'}{(\nu')^2} \text{Im} T_{\Delta\lambda}(\nu' + i\epsilon) \Big|_{\text{non-R}} \sim \frac{\alpha^2}{\nu_{\text{th}}} \sim 10^{-4} \text{ TeV}^{-2}, \quad (11)$$

where, for possible new physics states in the non-R loop in an underlying weakly interacting theory (if any), we expect the thresholds to be $\nu_{\text{th}} > (750 \text{ GeV})^2$. More precisely, in Quantum Electrodynamics (QED) with either a scalar or a spin- $\frac{1}{2}$ particle with charge $Q = 1$, the $\gamma\gamma$ cross section difference reaches a sharp global minimum with $\sigma_2(\nu) - \sigma_0(\nu) \gtrsim -8\alpha^2/\nu_{\text{th}}$ right after the production threshold due to the negative $\Delta\lambda = 0$ contribution, then a wider global maximum with $\sigma_2(\nu) - \sigma_0(\nu) \lesssim 2\alpha^2/\nu_{\text{th}}$ due to positive $\Delta\lambda = 2$ production, and finally a converging $1/\nu$ tail [13]. Thus, one finds that the pure QED one-loop amplitude for $\gamma\gamma \rightarrow \gamma\gamma$ fulfills the sum-rule (4) on its own and yields no correction [13]. Therefore, in order to get a contribution from these background loops to cancel the scalar resonance one in Eq. (10), one should incorporate effects beyond QED, which first enter at two loops. We find it very unlikely that these corrections are large enough to achieve this goal without entering a non-perturbative regime.

An elementary scalar that radiatively decays into two photons would not require in principle the addition of extra particles with $J_R \geq 2$, beyond new fermions with huge charges—or a huge number of components—if one considers a large diphoton partial width. However, as we have shown, achieving the latter and the required large background non-resonant contribution to the sum-rule implies a departure from perturbativity in the TeV range, where the BSM theory would enter a strongly coupled regime (see, e.g., Ref. [24]). Thus, one would expect to have composite states of any total angular momenta $J_R \geq 2$ lying in the non-perturbative energy range, as nothing forbids excitations with an arbitrary orbital momentum, similar to what one observes in QCD.

We are then left with the need of incorporating further resonance contributions with spin $J_R \geq 2$ to cancel out that from an $\mathcal{O}(\text{TeV})$ scalar candidate in Eq. (10).

Assuming that the sum-rule is dominated by the lowest spins and lightest resonances, the following relation for the lightest tensor resonance \mathcal{T} is obtained,

$$\Gamma_{\mathcal{T}\rightarrow\gamma\gamma} \approx \frac{\Gamma_{S\rightarrow\gamma\gamma}}{5} \left(\frac{M_{\mathcal{T}}}{M_S} \right)^3, \quad (12)$$

where we take the case in which the tensor decays only into $\gamma\gamma$ states with helicity difference $\Delta\lambda = 2$.⁴ Otherwise, we would obtain a larger value of $\Gamma_{\mathcal{T}\rightarrow\gamma\gamma}$ since the $\Gamma_{\mathcal{T}\rightarrow[\gamma\gamma]_0}$ would add up to the scalar contribution and a larger $\Gamma_{\mathcal{T}\rightarrow[\gamma\gamma]_2}$ would be needed to fulfill the sum-rule.

Finally, we note that these relations should eventually take into account additional resonances in the TeV region, correcting our lowest-resonance dominance assumption.

CONCLUSIONS

The late 2015 [1] and early 2016 [18] analyses by the ATLAS and CMS collaborations caused a huge excitement in our community, as they might have been the first direct evidence of a particle beyond the SM. On the contrary, mid 2016 data [2] did not show any significant diphoton excess. Of course, the diphoton channel continues to be one of the golden channels for new physics searches, which calls for a model-independent analysis of the consequences of a diphoton bump as we carried out in this paper.

The application of axiomatic quantum field theory to the diphoton production allows us to conclude that:

- i) A diphoton signal with an $\mathcal{O}(\text{TeV})$ mass confirmed as a broad spin-0 or 2 resonance will not come alone; as other resonances with different spin are needed to satisfy the derived sum rules. According to our discussion, this is obvious for the spin-0 case. Crossing symmetry implies that spin-2 (or higher) exchanges violate partial-wave unitarity, requiring the existence of a Regge-like tower with higher spin states. Odd-spins (3, 5, ...) need contributions from even spins—most likely $J_R = 0$ —to satisfy the sum rule (7) and require a tower of higher spins to restore partial-wave unitarity. All the reasoning in this article assumes a large width, of the order of the GeV; for much narrower resonances, the sum-rule cancellation can be achieved via non-R loop diagrams.
- ii) As a secondary conclusion, we remark that the proposed $\gamma\gamma$ sum-rules must be fulfilled for any value of the partial widths. Thus, in the case these had values much below the GeV, the tiny contribution from a resonance with $M_R \sim 1 \text{ TeV}$ to the sum-rule would still have to be cancelled out in a very precise way. However, in this case, it could be achieved by means of perturbative non-resonant BSM loops, which might be seen as a non-resonant background excess. Although it may be experimentally challenging, the direct evaluation of the sum rule from the data could give an estimate of the deviation from the SM and the presence of new physics.

⁴ This is the case, for instance, for massive gravitons G [7, 16, 17], the decay $G \rightarrow \gamma(\lambda)\gamma(\lambda')$ of which only occurs with $\Delta\lambda = 2$.

In the absence of precise enough data, the application of general quantum field theory principles allows the derivation of model-independent relations (applied here to the diphoton channel) that can be helpful in understanding potential new phenomena. For this reason, due to its simplicity and generality, the analysis shown in this letter is of high interest for the study of future diphoton searches at colliders.

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